

# Effective field theory analysis of composite higgsino-like and wino-like thermal relic dark matter

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**Abstract.** We study the effective field theory (including operators up to dimension five) of models in which dark matter is composite, consisting of either an electroweak doublet Dirac fermion ('higgsino-like dark matter') or an electroweak triplet Majorana fermion ('wino-like dark matter'). Some of the dimension-five operators in the former case cause mass splittings between the neutralino and chargino states, leading to a depleted rate of coannihilations and viable thermal relic dark matter with masses of the order of tens to hundreds of GeV rather than the usual pure higgsino thermal relic mass of 1 TeV. No such effects are found in the latter case (where the usual thermal relic mass is 3 TeV). Other operators, present for both wino- and higgsino-like dark matter, correspond to inelastic electromagnetic dipole moment interactions and annihilation through these can lead to viable models with dark matter masses up by an order of magnitude compared to the usual values.

**Keywords:** dark matter simulations, dark matter theory

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## 1 Introduction

The existence and identity of the dark matter in the universe have long been major issues in astronomy and particle physics [1]. There is now an overwhelming catalogue of astronomical evidence for the existence of dark matter, including from galactic rotation curves [2], large scale galaxy structure [3], the cosmic microwave background [4], and colliding galaxy clusters [5], but its identity remains a mystery.

A common and popular proposal to explain the particle nature of dark matter (assuming it is indeed a particle) is a fermion (Majorana or Dirac), which must of course be electrically neutral and an  $SU(3)$  color singlet, but may have some electroweak charge. The simplest possibilities are a Majorana singlet or triplet, or a doublet (which must be Dirac). Though we may generalise to larger multiplets [6], these examples have the nice feature that they correspond to the bino, wino, and higgsino of supersymmetric models. Even though supersymmetry will play no role in this work, we will freely employ its familiar nomenclature, referring, for example, to the electrically-neutral states in a multiplet as neutralinos, and to the charged states as charginos.

The models with a higgsino or a wino are particularly attractive because coannihilations [7–9], namely annihilation processes between the neutralinos and charginos, and the Sommerfeld enhancement [10] lead to natural thermal relic dark matter candidates close to the TeV scale. To wit, doublet models lead to a 1 TeV thermal relic while triplet models lead to a 3 TeV thermal relic [6, 11–17].

Unfortunately, the wino models are in conflict with the indirect detection constraints due to their strong tree-level annihilation into a pair of  $W$  bosons [18, 19]. The higgsinos evade these and the direct detection constraints (see e.g. [20–24]) as the cross sections for scattering from quarks are suppressed. In such models, even a tiny splitting of at least around 100 keV in the masses of the neutral and charged components also relax the direct detection limits as the scattering rates are now inelastic [25, 26] (for recent discussions of bounds in the SUSY case, see [27–30]). New mass terms which split the neutralino and chargino masses

also reduce the rate of coannihilations of dark matter, giving an increased relic density [7–9]. Such a tiny mass splitting can arise from new physics at much higher scales,<sup>1</sup> in which case it can be described in a completely model-independent fashion by dimension five operators in an effective field theory description [16, 40–42]. In other words, we regard dark matter as being composite.<sup>2</sup>

We thus see that there is a simple, viable description of dark matter, in terms of a Lagrangian containing only the Standard Model plus a higgsino, with terms of dimension up to and including five. Our goals in this work are, firstly, to explore the full parameter space of this model (with all possible dimension-five operators added), which we call higgsino-like dark matter and, secondly, to do the same for the corresponding model with the higgsino replaced by a wino, which we call wino-like dark matter.

There are two general classes of operators in each model. One class features a dimension-5 coupling between the dark matter and the Higgs field. Such operators can produce mass splittings between the neutral and charged components after the Higgs field is spontaneously broken [16, 42, 44]. The other class of operators correspond to electric and magnetic dipole moments for dark matter [45]. Such operators have been well-studied at the GeV scale [45–58] but here we analyse their effects near the TeV scale.

In section 2, we will introduce the formalism for each of the relevant dimension-5 operators. With these operators added, in section 3, we investigate the effect on the relic density and also compare to direct detection, indirect detection and collider bounds to test the viability of each of the relevant regions, while taking care to ensure that we remain within the region of validity of the effective field theory description. We show that higgsino-like models with a sizeable mass splitting can produce viable thermal relics down to a few hundred GeV, whilst direct annihilation through these operators can produce viable thermal relics at scales of up to tens of TeV. For wino-like models, the only viable thermal relics are at scales of up to tens of TeV.

## 2 Models

### 2.1 Doublet models (higgsinos)

Consider introducing a pair of Majorana spin-1/2 SU(2) doublets to the Standard Model. Though we borrow the notation and terminology of higgsinos from supersymmetric models, we seek to retain generality and so do not seek to specify the higher-scale model. The particles are represented by a pair of doublets of two-component spinors

$$\tilde{H}_u = \begin{pmatrix} \tilde{H}_u^+ \\ \tilde{H}_u^0 \end{pmatrix} \quad \text{and} \quad \tilde{H}_d = \begin{pmatrix} \tilde{H}_d^0 \\ \tilde{H}_d^- \end{pmatrix},$$

where we have used the round brackets for the representation of the gauge group.

The mass term in the Lagrangian is given by

$$\mathcal{L} \supset -\mu \epsilon^{ij} (\tilde{H}_u)_i (\tilde{H}_d)_j + \text{h.c.}, \quad (2.1)$$

<sup>1</sup>For example, in SUSY models, the splitting can arise due to radiative corrections from loops of heavy squarks [9, 31–33] or integrating out higher scale neutralinos [9, 34]. For other UV-complete models generating mass splittings, see [12, 35–39].

<sup>2</sup>For recent related work on composite dark matter, see [43].

for the antisymmetric tensor  $\epsilon^{ij}$  such that  $\epsilon^{12} = -\epsilon^{21} = +1$ , where the indices  $ij$  to indicate the gauge components. Note here that we can construct a pair of Dirac fermions with four-component spinors in the chiral basis

$$\tilde{H}^0 = \begin{bmatrix} \tilde{H}_u^0 \\ \tilde{H}_d^{0\dagger} \end{bmatrix} \quad \text{and} \quad \tilde{H}^+ = \begin{bmatrix} \tilde{H}_u^+ \\ \tilde{H}_d^{-\dagger} \end{bmatrix},$$

where the square brackets indicate spinor components. Hence, the mass term may also be written as

$$\mathcal{L} \supset -\mu \overline{\tilde{H}^0} \tilde{H}^0 + \mu \overline{\tilde{H}^+} \tilde{H}^+. \quad (2.2)$$

We can factor the Dirac spinors into a single SU(2) doublet

$$\tilde{H} = \begin{pmatrix} \tilde{H}^+ \\ \tilde{H}^0 \end{pmatrix}.$$

The model therefore also represents a single SU(2) doublet being added to the Standard Model. The mass is given by

$$\mathcal{L} \supset -\mu \overline{\tilde{H}_i} \tilde{H}_i. \quad (2.3)$$

The kinetic term for the Dirac doublet appears in the Lagrangian interacting under the SU(2)  $\times$  U(1) gauge symmetry of the Standard Model

$$\mathcal{L} \supset i \overline{\tilde{H}_i} \gamma^\mu \left( \partial_\mu - ig A_\mu^a \tau_{ij}^a - ig' Y B_\mu \delta_{ij} \right) \tilde{H}_j, \quad (2.4)$$

where  $A_\mu^a$  and  $B_\mu$  are, respectively, the SU(2) and U(1) Standard Model gauge bosons in the interaction basis,  $g$  and  $g'$  are the Standard Model gauge couplings,  $Y = +\frac{1}{2}$  is the weak hypercharge of the doublet, and  $\tau_{ij}^a$  are the generators of the SU(2) symmetry in the fundamental representation. Under the spontaneous symmetry breaking of the Higgs, we get the standard expressions of the mass-basis gauge bosons in the usual way:

$$W_\mu^\pm \equiv \frac{1}{\sqrt{2}} (A_\mu^1 \mp i A_\mu^2), \quad (2.5)$$

$$Z_\mu \equiv \frac{-g B_\mu + g A_\mu^3}{\sqrt{g^2 + g'^2}}, \quad (2.6)$$

$$A_\mu^\gamma \equiv \frac{g B_\mu + g' A_\mu^3}{\sqrt{g^2 + g'^2}}. \quad (2.7)$$

We have now recovered the usual interactions between an SU(2) fermionic doublet and the gauge bosons. The mass and covariant derivative terms are the only interactions allowed by the new doublets; there are no new gauge-invariant Yukawa terms in the theory as the only scalar particle in the model is the Higgs boson.

The only free parameter in this model is  $\mu$ , the mass parameter of the higgsinos. By imposing that the model must be a thermal relic dark matter candidate, then  $\mu$  is constrained to a mass of 1 TeV [12–14, 18, 27, 59, 60].

In our description, any additional physics above the scale of  $\mu$  is encoded in higher-dimensional effective field theory operators, which can result in modifications to both the masses and couplings of the higgsinos. Examples of such physics include the top-stop-higgsino

Yukawa from the MSSM, whose effects are dependent on the mass of the stop squark, or the addition of more weakly-interacting neutral fermions at higher mass scales (such as the bino or sneutrinos in the MSSM). The key advantage of following an effective field theory approach is that we do not need to specify the exact nature of the higher scale physics. In particular, supersymmetry is a possible, but not the only, UV completion of the theory.

In this work, we consider the eight independent dimension-5 operators for the pair of (Majorana) SU(2) doublet higgsinos. Four such operators correspond to an interaction with a scalar Higgs Boson  $H = \begin{pmatrix} H^+ \\ H^0 \end{pmatrix}$  as discussed in refs. [16, 42, 44], namely

$$\mathcal{O}_1^H = (H^\dagger)^i (\tilde{H}_u)_i (H^\dagger)^j (\tilde{H}_u)_j, \quad (2.8)$$

$$\mathcal{O}_2^H = \epsilon^{ij} \epsilon^{kl} (H)_i (\tilde{H}_d)_j (H)_k (\tilde{H}_d)_l, \quad (2.9)$$

$$\mathcal{O}_3^H = \epsilon^{jk} (H^\dagger)^i (\tilde{H}_u)_i (H)_j (\tilde{H}_d)_k, \text{ and} \quad (2.10)$$

$$\mathcal{O}_4^H = \epsilon^{jk} (H^\dagger)^i (\tilde{H}_d)_i (H)_j (\tilde{H}_u)_k. \quad (2.11)$$

As noted in ref. [42], no terms of the form  $\epsilon^{ij} \epsilon^{kl} (H)_i (H)_j (\tilde{H}_d)_k (\tilde{H}_d)_l$  are permitted due to the fact that the  $H$  bosons are symmetric. Given the two-dimensional SU(2) generators  $\tau_{ij}^a$ , it is possible to write these operators in the form  $(H_i^\dagger \tau_{ij}^a \tilde{H}_j)(H_k^\dagger \tau_{kl}^a \tilde{H}_l)$  using the Fierz relation of the generators  $\tau_{ij}^a \tau_{kl}^a = 2\delta_{il}\delta_{jk} - \delta_{ij}\delta_{kl}$ .

Terms like these arise from, for example, integrating the gauginos out from the MSSM. They naturally allow for modifications to the higgsino masses after the Higgs boson obtains a vacuum expectation value. Physically, the mass modification represents a small mixing of the masses to a significantly heavier state.

In addition, we introduce four operators which together contain the electric and magnetic dipole moments of the higgsinos, namely

$$\mathcal{O}_5^H = \widetilde{H}_i \sigma^{\mu\nu} B_{\mu\nu} \tilde{H}_i, \quad (2.12)$$

$$\mathcal{O}_6^H = \widetilde{H}_i \tau_{ij}^a \sigma^{\mu\nu} W_{\mu\nu}^a \tilde{H}_j, \quad (2.13)$$

$$\mathcal{O}_7^H = \widetilde{H}_i \sigma^{\mu\nu} B_{\mu\nu}^* \tilde{H}_i, \text{ and} \quad (2.14)$$

$$\mathcal{O}_8^H = \widetilde{H}_i \tau_{ij}^a \sigma^{\mu\nu} (\widetilde{W}^a)^*_{\mu\nu} \tilde{H}_j, \quad (2.15)$$

where  $B_{\mu\nu}$  and  $W_{\mu\nu}^a$  are the field strength tensors for the U(1) and SU(2) gauge fields respectively

$$B_{\mu\nu} = \partial_\mu B_\nu - \partial_\nu B_\mu, \quad (2.16)$$

$$W_{\mu\nu}^a = \partial_\mu A_\nu^a - \partial_\nu A_\mu^a + f^{abc} A_\mu^b A_\nu^c, \quad (2.17)$$

and  $B_{\mu\nu}^*$  and  $(W^a)^*_{\mu\nu}$  are the dual field strength tensors  $B_{\mu\nu}^* = \epsilon_{\mu\nu\sigma\rho} B^{\sigma\rho}$  and  $(W^a)^*_{\mu\nu} = \epsilon_{\mu\nu\sigma\rho} (W^a)^{\sigma\rho}$ . Note here that the  $\tau_{ij}^a$  term in eqs. (2.13) and (2.15) mixes the states of the doublet. As the neutral states are mixed, then if there is a mass splitting the resultant dipole moments will be inelastic [25].

The collection of operators in eqs. (2.8)–(2.15) represent the complete set of gauge-invariant operators at dimension 5 up to the usual redundancies in effective field theory, as dimensional analysis and symmetry restrict us to operators involving two fermions and two bosons.

The neutral components of operators  $\mathcal{O}_5^H$  and  $\mathcal{O}_6^H$  correspond to an electroweak magnetic dipole operator, and the neutral components of operators  $\mathcal{O}_7^H$  and  $\mathcal{O}_8^H$  correspond to an electroweak electric dipole operator, as can be seen using the identity  $\frac{i}{2}\sigma^{\sigma\rho}\epsilon_{\mu\nu\rho\sigma} = \sigma^{\mu\nu}\gamma^5$ .<sup>3</sup> Electric and magnetic dipolar dark matter have been studied extensively in the literature for electromagnetically interacting U(1) models at masses of a few to a hundred GeV [46–58]. Our treatment here considers an electroweak SU(2) × U(1) model, which also allows interactions between the charged higgsinos and  $W^\pm$  bosons. Also, note that the electric dipole-like terms with the dual field strength tensors are  $CP$ -violating terms.

To examine the properties of these operators, we collate the dimension-5 terms into the effective Lagrangian

$$\mathcal{L} \supset \sum_{i=1}^8 \frac{c_i}{\Lambda} \mathcal{O}_i^H + \text{h.c.}, \quad (2.18)$$

for some coupling coefficients  $c_i$  and UV-cutoff scale  $\Lambda$ . In principle, some of the  $c_i$  may be complex, but here we will assume them to be all real.

As eqs. (2.8)–(2.11) give four operators that provide a coupling to the Higgs boson, these operators will provide additional mass terms to  $\tilde{H}$  after electroweak symmetry breaking [42]. Notably,  $\mathcal{O}_1^H$  and  $\mathcal{O}_2^H$  create a splitting of the masses of the two neutral states by placing terms on the main diagonal of an otherwise anti-diagonal mass matrix. With these terms switched on, the magnitude of the eigenvalues of the mass matrix are no longer equal.  $\mathcal{O}_3^H$  adds off-diagonal terms, because only the neutral component of the Higgs gets a vev. Finally,  $\mathcal{O}_4^H$  shifts the mass of the charged state. Hence we obtain the mixing matrix

$$\mathcal{L} \supset \tilde{H}_i^0 \mathcal{M}_{ij} \tilde{H}_j^0 = -\frac{1}{2} \begin{pmatrix} \tilde{H}_u & \tilde{H}_d \end{pmatrix} \begin{pmatrix} \frac{v^2 c_1}{\Lambda} & -\mu + \frac{v^2 c_3}{2\Lambda} \\ -\mu + \frac{v^2 c_3}{2\Lambda} & \frac{v^2 c_2}{\Lambda} \end{pmatrix} \begin{pmatrix} \tilde{H}_u \\ \tilde{H}_d \end{pmatrix}. \quad (2.19)$$

The mass matrix is diagonalised for complex parameters  $\mu$ ,  $c_1$ ,  $c_2$  and  $c_3$  through the usual diagonalization of complex symmetric matrices [63]; the masses obtained are given by

$$m_{1,2}^2 = \frac{1}{2} \left( 2|\tilde{\mu}|^2 + \frac{v^4(|c_1|^2 + |c_2|^2)}{\Lambda^2} \pm \sqrt{\frac{v^8(|c_1|^2 - |c_2|^2)^2}{\Lambda^4} + 4\frac{v^4|c_2\tilde{\mu}^* + c_1^*\tilde{\mu}|^2}{\Lambda^2}} \right). \quad (2.20)$$

In the limit that the parameters are real (more specifically, that the complex phases are equal), the neutral masses are given by

$$m_1 = \left| \tilde{\mu} - \frac{c_1 + c_2}{2\Lambda} v^2 \right| \quad (2.21)$$

and

$$m_2 = \left| \tilde{\mu} + \frac{c_1 + c_2}{2\Lambda} v^2 \right|, \quad (2.22)$$

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<sup>3</sup>It is well known that Majorana particles cannot possess electric or magnetic dipole moments; the first permissible electromagnetic term is the dimension-6 anapole moment [61]. The electric and magnetic dipoles are both  $CPT$ -odd, so for a particle that is  $CPT$  self-conjugate, the terms cannot exist in a  $CPT$  consistent theory [62]. However, the restriction does not apply in this theory because of the presence of two SU(2) doublets being introduced simultaneously. Where both of the doublets have identical masses, the accounting trick of making the 2 doublets into one Dirac fermion in eq. (2.3) resolves the issue. Where the masses are split via operators  $\mathcal{O}_1^H$  and  $\mathcal{O}_2^H$ , these equations describe an inelastic dipole between the two mass eigenstates, or between the neutralino eigenstates and the chargino eigenstates.

with the charged mass given by

$$m_{\pm} = \left| \mu + \frac{c_4}{2\Lambda} v^2 \right|. \quad (2.23)$$

Here we have defined

$$\tilde{\mu} = \frac{1}{2} \sqrt{\frac{(c_1 - c_2)^2}{\Lambda^2} v^4 + 4\bar{\mu}^2}, \quad (2.24)$$

and

$$\bar{\mu} = \mu - \frac{c_3}{2\Lambda} v^2, \quad (2.25)$$

where  $v \simeq 246$  GeV is the Higgs vacuum expectation value. Note that in the limit where  $c_1 \simeq c_2$  as we consider later in this work, then  $\tilde{\mu} \simeq \bar{\mu}$ . Finally, the transformation from the interaction basis to the mass basis is given by

$$\begin{pmatrix} \tilde{\chi}_1^0 \\ \tilde{\chi}_2^0 \end{pmatrix} = \begin{pmatrix} \cos \theta & -\sin \theta \\ \sin \theta & \cos \theta \end{pmatrix} \begin{pmatrix} \tilde{H}_d^0 \\ \tilde{H}_u^0 \end{pmatrix}, \quad (2.26)$$

with

$$\tan \theta \simeq 1 + \frac{c_2 - c_1}{2\mu\Lambda} v^2. \quad (2.27)$$

It must be noted here that eqs. (2.21) and (2.22) imply that, in the case where  $\tilde{\mu} \sim \frac{v^2}{\Lambda}$ , the mass of the lighter particle approaches zero, while the heavier of the two neutralinos and the chargino remain of  $\mathcal{O}(\mu)$ . Thus a large mass difference between the two masses can be generated, which is crucial for the effect on the relic density. Where there is a sizeable mass difference between the chargino and neutralino, the rate of coannihilations is significantly reduced, greatly impacting the thermal relic density. This region is capable of producing feasible thermal relic dark matter with significantly altered dark matter mass.

## 2.2 Triplet models (winos)

Next, we introduce a weakly-interacting, Majorana, SU(2)-triplet,  $\widetilde{W}$ , with zero weak hypercharge. The model is analogous to a supersymmetric wino, and again we borrow the notation and terminology, but stress that such a model need not be supersymmetric. The components are represented in the usual way for winos, with two charged components  $\widetilde{W}^{\pm}$  and a neutral component  $\widetilde{W}^0$ , which may be written in doublet form as  $\widetilde{W}_i = (\widetilde{W}^+, \widetilde{W}^0, \widetilde{W}^-)$ . The mass term is given by

$$\mathcal{L} \supset -\frac{1}{2} M_2 \widetilde{W}_i \widetilde{W}_i. \quad (2.28)$$

Since there is zero hypercharge, there is only an interaction term corresponding to the SU(2) symmetry of the Standard Model

$$\mathcal{L} \supset i \widetilde{W}_i \bar{\sigma}^{\mu} \left( \delta_{ij} \partial_{\mu} - ig A_{\mu}^a \frac{T_{ij}^a}{2} \right) \widetilde{W}_j + \text{h.c.}, \quad (2.29)$$

where  $T_{ij}^a$  are the three-dimensional representations of the generators of SU(2) and  $\sigma^{\mu}$  is the . These dimension-4 interactions can produce a thermal relic particle with mass 3 TeV [11–15, 18]. Variations to the mass of the relic occur if the neutral component of the triplet is

mixed with a singlet or doublet (i.e. a bino-like or higgsino-like particle). As for the higgsino, we seek to model the effect of any such (heavy) additional new physics through effective field theory operators.

Dimensional analysis and symmetry restrict us at dimension 5 to a coupling between two fermionic states and two bosonic states. As for the doublet case, we expect a coupling to the Higgs bosons of the form

$$\mathcal{O}_1^W = (H^\dagger)^i (H)_i \widetilde{W}^2. \quad (2.30)$$

The operator in eq. (2.30) contains two Higgs bosons, and hence we expect to see shifts or splittings in the mass spectrum after electroweak symmetry breaking. Unlike the doublets, where the masses occur on the anti-diagonal of the mixing matrix, here we have a direct contribution to the doublet masses

$$\mathcal{L} \supset \widetilde{W}^i \left( M_2 + \frac{d_1 v^2}{\Lambda} \right) \widetilde{W}^i. \quad (2.31)$$

Here, the chargino and neutralino masses are equal regardless of the scale of the dimension-5 coupling.

For the interaction between the triplet fermion and the gauge bosons at dimension 5, it is possible to write down two gauge-invariant terms, connecting the  $\widetilde{W}$  fields either to the Abelian  $B_{\mu\nu}$  field strength tensor, and the non-Abelian  $W_{\mu\nu}$  field strength tensor. Each of these terms also has an equivalent term with the dual field strength tensor which is  $CP$  violating.

The two terms are

$$\mathcal{O}_2^W = \widetilde{W}_i T_{ij}^a \sigma^{\mu\nu} W_{\mu\nu}^a \widetilde{W}_j, \text{ and} \quad (2.32)$$

$$\mathcal{O}_3^W = \widetilde{W}_i T_{ij}^a \sigma^{\mu\nu} (W_{\mu\nu}^a)^* \widetilde{W}_j. \quad (2.33)$$

We cannot introduce operators containing  $\overline{W}_i \sigma^{\mu\nu} B_{\mu\nu} \overline{W}_i$  or  $\overline{W}_i \sigma^{\mu\nu} B_{\mu\nu}^* \overline{W}_i$  as though they are Lorentz invariant, they vanish identically via the antisymmetry of the fermionic components. Indeed, these components would have corresponded to an electric dipole operator which is forbidden for Majorana particles under a  $CPT$ -invariant theory [61, 62].

However, by introducing a gauge field with matrix components, it is possible to form a Lorentz invariant via a coupling between the antisymmetric components of  $T_{ij}^a$  and the antisymmetric combination  $\overline{W}_i \overline{W}_j$ . Hence the terms in eqs. (2.10) and (2.11) can avoid the constraints on magnetic dipoles in  $CPT$ -invariant theories.

Finally, the Lagrangian is compiled as

$$\mathcal{L} \supset \sum_{i=1}^3 \frac{d_i}{\Lambda} \mathcal{O}_i^W + \text{h.c.}, \quad (2.34)$$

for coupling parameters  $d_i$  and UV-cutoff  $\Lambda$ . As usual in EFT, we take the coefficients  $c_i$  to be order one or smaller. In principle some may be complex, but here we take them all to be real.

We have now introduced all dimension-5 terms of the higgsinos and winos introduced separately into the theory. These terms account quite generally for the leading effects of physics at higher scales. The new terms provide both electric and magnetic dipole-like interactions, and couplings to the Higgs provide mass terms. There is a regime where if the effective field theory couplings is sufficiently large, the masses may split, cancel or dominate, which is relevant to the coannihilation rates in the relic density calculations.



### 3 Relic density

Any reasonable dark matter model must account for the abundance of dark matter in the universe. In the thermal relic dark matter scenario, the dark matter is in thermal equilibrium with the Standard Model in the early universe. At the freeze-out epoch, the expansion of the universe prevents dark matter from annihilating, and the remaining relic dark matter density remains constant through to the present day. The thermal relic density is dependent on the size of the annihilation cross-section. Although there are models which explain the abundance of dark matter through non-thermal mechanisms [64], in this work we focus solely on thermal models.

In higgsino-like or wino-like dark matter, the only free parameter affecting the annihilation cross section at the renormalizable level is the mass parameter  $\mu$  or  $M_2$ , respectively. Here, the only relevant interactions that can contribute to the dark matter annihilation in the early universe are with the gauge bosons. In the Lagrangian, these interactions arise in the covariant derivative, and lead to, for example,  $s$ -channel annihilation via a  $Z$  boson or  $t$ -channel annihilation to a pair of  $W$  bosons. The coupling strength is therefore given by the electroweak coupling parameters  $g$  and  $g'$ , which are fixed in the Standard Model.

As the annihilation cross section increases, the relic density decreases as the dark matter abundance is depleted more efficiently. Hence, aside from resonance effects near the masses of the gauge bosons, there is a direct proportionality between the relic density and the dark matter mass parameters. The relic density measurement then uniquely constrains the relevant mass parameter. For higgsino-like dark matter at dimension 4 and below, the relic density constrains the mass to around 1 TeV, whereas wino-like dark matter at dimension 4 and below is constrained to around 3 TeV. Searches for neutralino dark matter have focussed on these mass regions, with the 3 TeV thermal relic wino dark matter being ruled out by indirect detection experiments [18, 19], with the 1 TeV thermal higgsino requiring mass splittings, albeit small ones, in order not to be ruled out by direct detection experiments.

In this work, we investigate the effect of introducing the eight additional dimension-5 operators for higgsino-like dark matter in eqs. (2.8)–(2.15) or the five dimension-5 operators for wino-like dark matter in eqs. (2.30), (2.32) and (2.33) to the relic density of the respective particles. In particular, we consider the effect on the relic density spectrum owing to the modifications to the higgsino-like masses in eqs. (2.21) and (2.22). We calculate the relic densities using the `micrOMEGAs` computer code [65–67], using the inbuilt `CalcHEP` [68] and `LanHEP` [69] functionality to generate the Feynman diagrams and associated squared matrix elements. Sommerfeld enhancements are then approximated for each model using the procedure in ref. [70].

For each operator we consider, we vary two parameters that contribute to the mass spectrum of the higgsino-like or wino-like particle. These are the mass parameter  $\mu$  or  $M_{1,2}$  and the UV cut-off associated with the dimension-5 coupling  $\Lambda$ . In all cases, we set  $c_i, d_i = \pm 1$ , which represents roughly the largest possible allowed coupling magnitude. For each of the operators, there is a region of the parameter space where the masses are greater than the cut-off scale. In these regions, the effective field theory is not valid, in that the details of the higher-scale physics become relevant; we exclude such regions from the results.

For each of our models, we consider the relic density and constraints from experimental results. We use compare the simulated relic density to the observed relic density from Planck of  $\Omega h^2 = 0.1188 \pm 0.0010$  [71]. The regions which provide the correct relic density are shown in each of the figures as a thick green line. A further constraint arises from collider

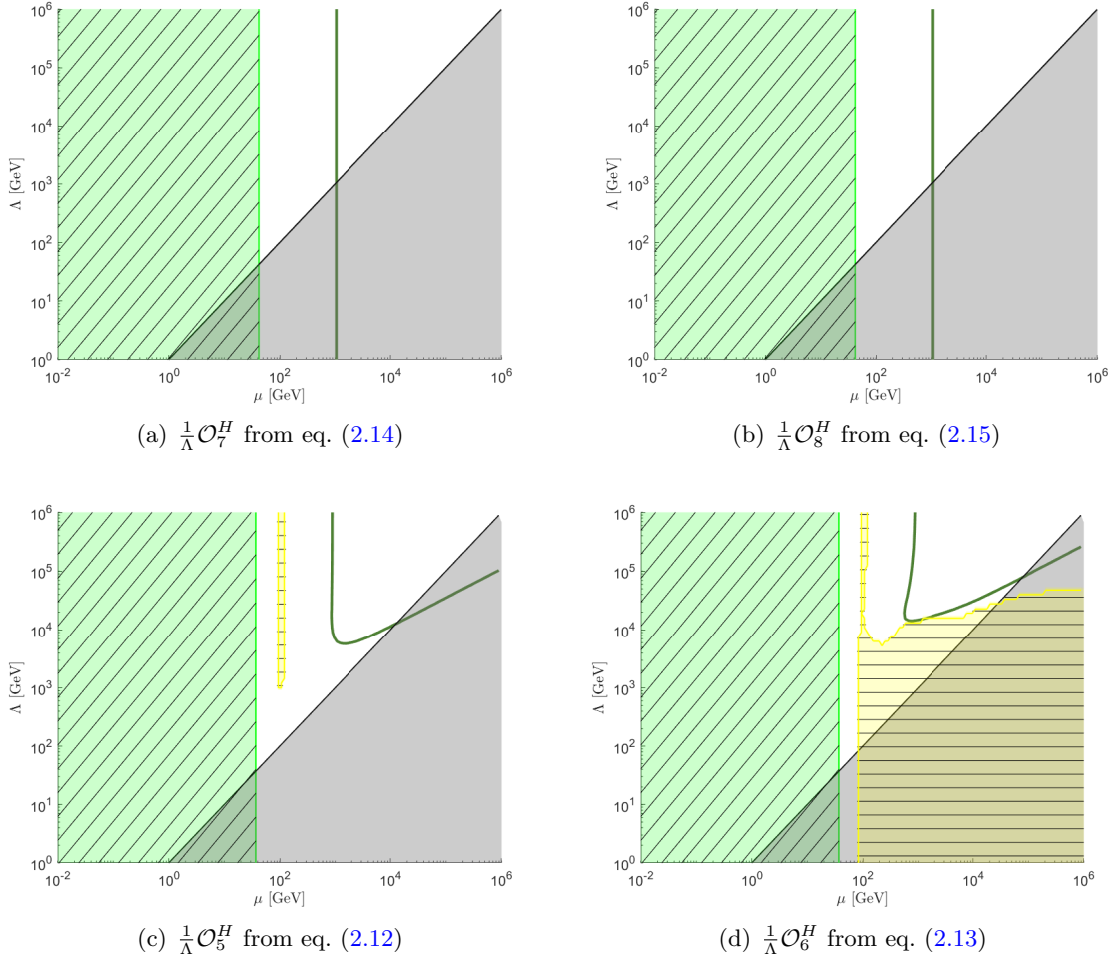
experiments, notably the invisible width of  $Z$  boson decays as measured by LEP [72–77]. The mass of the charged component of the higgsino-like or wino-like fermions must be greater than 103.5 GeV [76, 78] to avoid a contribution to  $\Gamma_Z$ . However, no such constraint can be placed on the neutral component of the higgsino-like fermion, as there is no direct coupling to the  $Z$  boson.

There are also indirect detection constraints imposed by astronomical observations. Searches for dark matter using gamma ray astronomy at the TeV scale include satellite-based (e.g. Fermi-LAT) and ground-based (e.g. HESS) telescopes looking for the by-products of dark matter annihilation in both the galactic centre [79] and extra-galactic sources [80]. In this work we impose the Fermi-LAT dwarf spheroidal [80] and HESS galactic centre constraints [81]; the former bounds are stronger for masses of around 100 GeV, the latter is stronger for masses above 1 TeV. Before the introduction of dimension-5 operators, wino-like thermal dark matter is in tension with the Fermi-LAT constraints [18], subject to standard assumptions on the nature of the galactic halo, especially in the galactic centre. The strongest bound for the present work is in the  $W^+W^-$  channel, as it is directly produced by the electric and magnetic dipole interactions in eqs. (2.12)–(2.15) and eqs. (2.30), (2.32) and (2.33). We calculate the annihilation rates for each point in our parameter space using the `micrOMEGAs` code [65–67], followed by a calculation of the Sommerfeld enhancement [70, 82, 83]. All operators are affected by the constraints on the  $W^+W^-$  annihilation channel when the mass of the LSP is near the resonance at the  $W$  boson mass, regardless of the dimension-5 couplings.

The direct detection bounds on the higgsino-like models are weak at tree level: the  $t$ -channel scattering is suppressed as the two neutral components are of near equal mass creating a cancellation in the vertex with the  $Z$  boson and there is no  $s$ -channel annihilation via a heavy squark as there is in supersymmetric models [84, 85]. The result is an expected scattering cross section below the neutrino floor for direct detection. The dimension-5 operators which do have a coupling to the  $Z$  boson also undergo similar cancellations at tree level. Meanwhile, for the wino-like models there are no diagrams at tree level that permit elastic scattering of the dark matter off nucleons. In general, direct detection bounds are not strong enough to rule out any significant regions of the parameter space of the higgsino-like particles, at least under the usual conditions where there is a tiny mass splitting between the higgsino and the chargino [25, 26]. The exception occurs for regions where the LSP is the chargino, that is, when  $\frac{c_1+c_2}{2\Lambda}v^2 > \mu + \frac{c_4}{2\Lambda}v^2$ . For the higgsino, this corresponds to small  $\mu$  and a small cut-off scale. Such regions are strongly excluded as they represent charged dark matter, which has been excluded by numerous experiments, for example, the abundance of superheavy isotopes of hydrogen in sea water [86], and are shown in the figures in this section as the direct detection exclusion zones. Although we also compute the direct detection rates for the XENON1T [22] results, none of the bounds are strong enough to be shown in our figures.

### 3.1 Higgsino electric dipole interactions

The parameter space for the electric dipole operators in eqs. (2.14)–(2.15) are shown in figures 1(a)–(1(b)). Both of these operators are CP-violating. Perhaps unsurprisingly, the effect of these CP-violating operators on the relic density and other experimental constraints is negligible. Rather, all of the relevant physics is determined by the dimension-4 terms in the Lagrangian, namely the interaction in eq. (2.4). Figures 1(a)–1(b) thus essentially represent the null hypothesis, where there is no impact from the higher dimension operators, which may be contrasted with the effects of the other operators. Hence, we may highlight the features



**Figure 1.** Parameter space for electric (top) and magnetic (bottom) dipole operators. The horizontal axes show the mass parameter  $\mu$  and the vertical axes UV-cutoff allowed for the parameter point, determined with the numerical coefficients  $c_i = 1$ . The contours of parameter values which satisfy the relic abundance are shown as a solid green line. The regions where the EFT is not valid are shaded grey. The regions excluded by indirect detection are shaded yellow with horizontal hatching. The regions excluded by LEP are shaded green with upwards diagonal hatching.

of the parameter space before the inclusion of more interesting dimension-5 operators. Note that for these operators, the mass of the neutralinos and chargino is given by  $\mu$ . The lower right of the parameter space (small  $\Lambda$ , large  $\mu$ ) represents the region where the effective field theory is not valid, as the physics occurs above the UV-cutoff. The region is bounded by the line  $\mu = \Lambda$ , which is the largest possible mass with a valid EFT.

The observed relic density occurs at  $\mu \simeq 1$  TeV. For larger masses, the relic density increases and the universe becomes over-closed without some additional non-thermal mechanism to reduce the abundance. For smaller masses, the relic density decreases, apart from a resonance near the gauge boson masses. Also near the gauge boson masses, the resonance annihilation leads to indirect detection constraints. Lower masses are then subject to the LEP constraints on both the neutralino and chargino masses.

### 3.2 Higgsino magnetic dipole interaction via $W$

Figures 1(c)–1(d) show the parameter space for the operators in eqs. (2.12) and (2.13) respectively. Together, these operators represent the magnetic dipole interaction of the higgsino. These operators provide point annihilations directly from two higgsinos to two gauge bosons as well as derivative interactions to a single gauge boson. The interactions increase in strength for decreasing  $\Lambda$ . The annihilation processes to a region of the parameter space for  $\Lambda < 10$  TeV and  $\mu > m_W$  where the magnetic dipole interactions dominate the regular annihilation processes. The increase in the annihilation cross section in this region depletes the otherwise over-abundant dark matter, allowing for the relic density to be satisfied for masses greater than 1 TeV. The region allows for thermal relic higgsinos to have masses up to around 70 TeV, which corresponds to a magnetic dipole on the order of  $10^{-3}$  to  $10^{-4}$  times the proton magnetic moment.

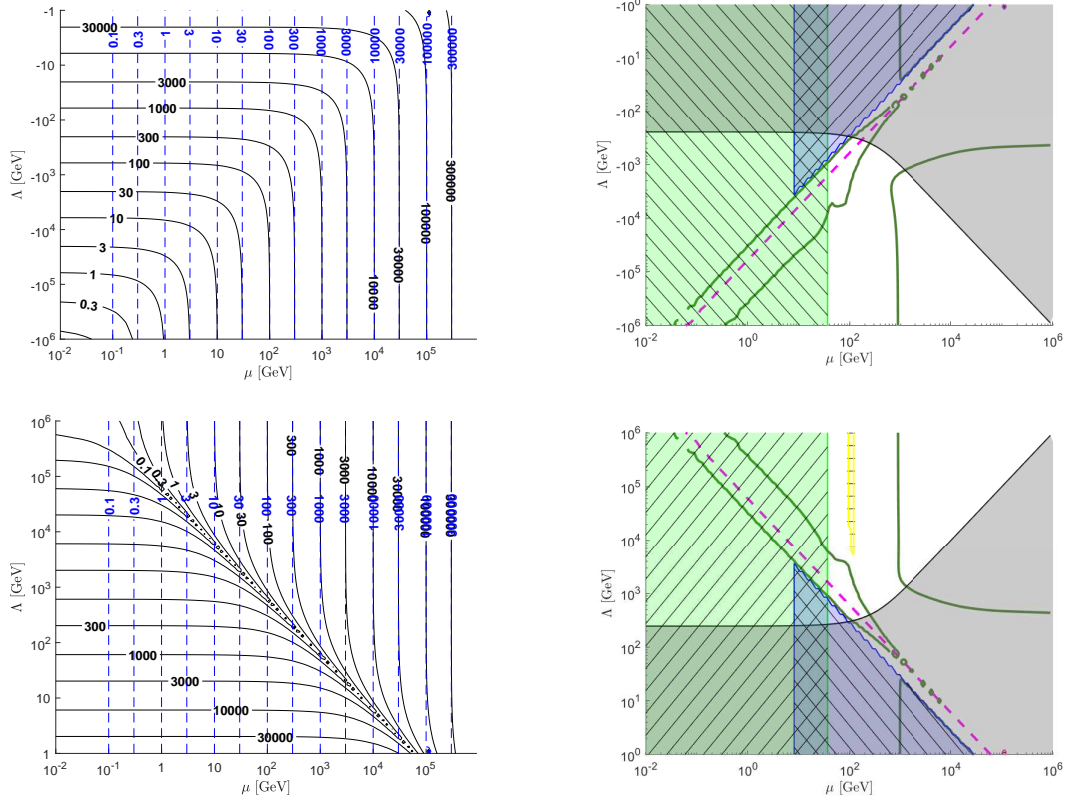
The operators also can be seen through indirect detection experiments, which can look directly for the annihilation products from the magnetic dipole. The operator  $\mathcal{O}_6^H$  allows for direct annihilation into a  $W^+W^-$  pair, whereas the  $\mathcal{O}_5^H$  operator can only annihilate into a combination of  $Z$  bosons and photons. The indirect detection bounds are given for the  $W^+W^-$  channel only [80]. Meanwhile, the reduced abundances due to the additional annihilation reduce the total amount of annihilation in the  $W^+W^-$  channel for indirect detection searches, weakening the bounds.

### 3.3 Higgs-higgsino interaction with neutralino mass splitting

The Higgs-higgsino interaction terms provide significantly different phenomena than the higgsino dipole interactions. Most notably, the dimension-5 terms will contribute mass terms for the higgsino after the Higgs Boson undergoes symmetry breaking. The mass of the two higgsinos and chargino become non-degenerate. Figures 2(a) and 2(b) show the parameter space for the Lagrangian terms  $\frac{c_{1,2}}{\Lambda}(\mathcal{O}_1^H + \mathcal{O}_2^H)$  from eqs. (2.8)–(2.9) where, for simplicity, we have assumed  $c_1 = c_2 = c_{1,2} = \pm 1$ . These terms are combined as they provide for a mass splitting between the two neutralinos through eqs. (2.21)–(2.22). The consequence of the two terms being equal is that we do not have an offset to  $\bar{\mu}$  per eq. (2.24), which is dependent on the difference between the two coefficients. The mass of the chargino does not depend on  $c_1$  or  $c_2$ , and so is given by  $\mu$ .

For  $\mu \gg \frac{v^2}{\Lambda}$ , that is,  $\mu$  and  $\Lambda$  are both large, all of the masses are degenerate and equal to  $\mu$ . For  $\mu \ll \frac{v^2}{\Lambda}$ , that is,  $\mu$  and  $\Lambda$  are both small, the higgsino masses are both given by  $\frac{c_{1,2}v^2}{\Lambda}$ , but the chargino mass is still  $\mu$  and there is a sign difference between the two neutralinos. There is therefore a mass splitting between the chargino and the higgsinos. Notably, in this region, the chargino becomes the lightest supersymmetric particle, instead of the higgsino. The region is strongly ruled out as it interacts too strongly with the Standard Model through the presence of the electric charge.

In the regime that  $\mu \sim \frac{v^2}{\Lambda}$ , shown in figures 2(b) as a dashed magenta line, the mass of the lightest higgsino approaches zero, while the mass of the chargino remains at  $\mu$ . The heaviest higgsino is now twice as massive as the chargino. Hence, the degeneracy between the chargino and the neutralino is broken. There is a significant impact on the relic density as the dominant annihilation mode is the coannihilation mode where near-degenerate higgsinos and charginos annihilate into fermions via a  $W$  boson. Indeed, in the limit that the winos and binos (if present) have heavy mass, that is, the two neutralinos are an identical mix of the two higgsinos, the annihilation between neutral higgsinos vanishes [87]. This is because



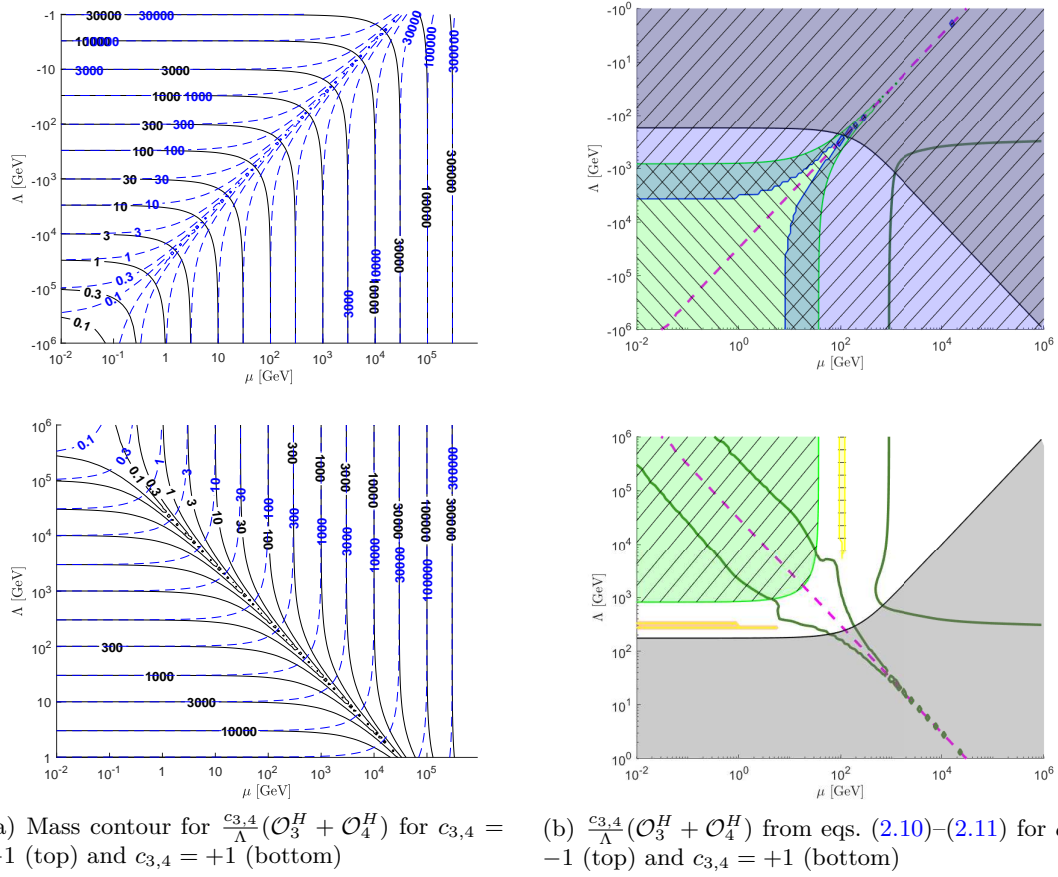
(a) Mass contour for  $\frac{c_{1,2}}{\Lambda}(\mathcal{O}_1^H + \mathcal{O}_2^H)$  with  $c_{1,2} = -1$  (top) and  $c_{1,2} = +1$  (bottom)      (b)  $\frac{c_{1,2}}{\Lambda}(\mathcal{O}_1^H + \mathcal{O}_2^H)$  from eqs. (2.8)–(2.9) with  $c_{1,2} = -1$  (top) and  $c_{1,2} = +1$  (bottom)

**Figure 2.** Mass contour plots (left) and parameter space (right) for Higgs-higgsino operators. Mass contour shows the lightest neutralino mass in black and the lightest chargino mass in blue as a function of the input parameters. The top graphs show negative values of the coupling constant and the bottom graphs show the positive values. The region where  $\mu \simeq \frac{v^2}{\Lambda}$  is marked as a dashed magenta line. The regions excluded by direct detection are shaded blue with downwards diagonal hatching. For a description of the parameter space plots, see the caption of figure 1.

of a cancellation between the two higgsino-higgsino- $Z$  boson terms in the Lagrangian. In the canonical case, the annihilation via the chargino and  $W$ -boson exchange is dominant [8]. Hence, by removing the mass degeneracy we no longer have a viable annihilation processes, leaving the universe to be overclosed by higgsino dark matter. There is then a continuous variation away from the case where  $\mu \sim \frac{v^2}{\Lambda}$  where the coannihilation pathway becomes unlocked. Along this variation is a point whereby the annihilation rate corresponds to the relic density. The dark matter mass along this contour ranges from  $m_\chi \sim 50$  GeV to  $m_\chi \sim 100$  GeV. The lower bound arises from the LEP bounds, whereas the upper bound is limited by the range of the EFT.

### 3.4 Higgs-higgsino interaction without neutralino mass splitting

Figures 3(a) and 3(b) show the mass contour and parameter space for the operators which do not split the neutralino masses  $\frac{c_{3,4}}{\Lambda}(\mathcal{O}_3^H + \mathcal{O}_4^H)$  from eqs. (2.10)–(2.11). The parameter space is similar to the mass splitting operators above, however the two neutralinos are still degenerate



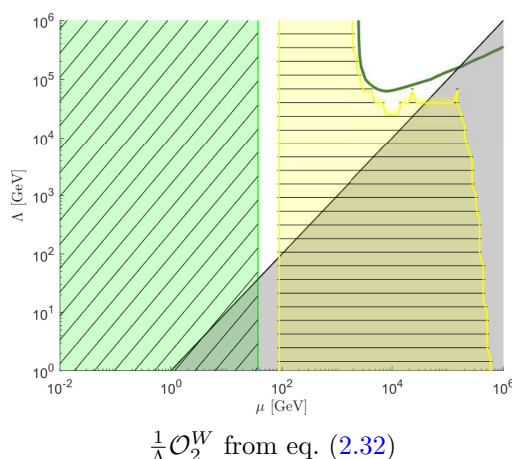
**Figure 3.** Mass contour plots (left) and parameter space (right) for Higgs-higgsino operators. Mass contour shows the lightest neutralino mass in black and the lightest chargino mass in blue as a function of the input parameters. The regions excluded by direct detection are shaded blue with downwards diagonal hatching. For a description of the parameter space plots, see the caption of figure 1.

throughout as the  $c_3$  amounts to a shift in the  $\mu$  components in the mass matrix. Now, the chargino mass is also affected by the  $c_4$  component in eq. (2.23), which is disconnected from the modified  $\mu$  parameter in eq. (2.25).

For  $\mu \gg \frac{v^2}{2\Lambda}$ , that is, both  $\mu$  and  $\Lambda$  are large, the masses of both the higgsinos and charginos are degenerate and equal to  $\mu$ . There is no impact from the dimension 5 operators in this regime. For  $\mu \ll \frac{v^2}{2\Lambda}$ , that is, both  $\mu$  and  $\Lambda$  are small, the masses of the higgsinos are again both large and degenerate, with  $m_1 = m_2 = \frac{v^2 c_3}{2\Lambda}$  and  $m_{\pm} = \frac{v^2 c_4}{2\Lambda}$ . As we have set  $c_3 = c_4$ , the chargino and neutralino are also still degenerate.

Note that the masses in the parameter space are symmetric about  $\mu = \frac{v^2}{2\Lambda}$ . The indirect detection bounds and relic density bounds are therefore also reflected about the line. Along the line  $\mu \sim \frac{v^2}{\Lambda}$ , shown as a dashed magenta line in figure 3(b), the masses of neutralinos approach zero but the chargino mass is doubled, as per the mass splitting case. Again, the annihilation processes that contribute the most to the higgsino relic density are the coannihilations with the chargino, for example the  $t$ -channel annihilation to  $W$ -bosons. Where there is the near cancellation between the two terms, the coannihilation processes are no longer efficient and so the annihilation does not proceed. Like the mass splitting case, there





**Figure 4.** Parameter space for wino magnetic dipole operator. For a description, see the caption of figure 1.

is a continuous variation between the highly fine-tuned cancellation and the standard case, along which there is a set of parameters where the relic density is met. There is a viable region for relic higgsino dark matter with masses  $m_\chi = 80$  GeV to  $m_\chi = 120$  GeV. The mass cancellation also affects the relic density near 1 TeV, reducing the allowed region as low as  $m_\chi = 500$  GeV. Finally, below  $\Lambda = 10^3$  GeV, the direct annihilation from the additional terms becomes significant and dominates the annihilations, but this only corresponds to small region of the viable parameter space, because for  $\Lambda < v$  the effective field theory breaks down.

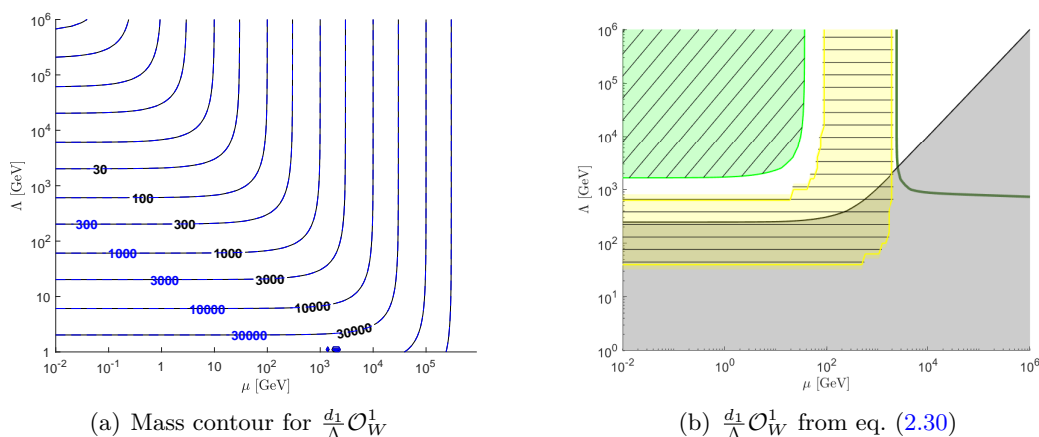
However, when the coupling parameter is negative, the mass cancellation now occurs in the chargino rather than the neutralino as per eq. (2.23). In the valley where  $\mu \sim \frac{v^2}{\Lambda}$ , there is no additional regime with a viable neutralino relic density, and the parameter space is excluded since the chargino becomes the lightest particle in the model. As the argument of the complex coupling constant transitions from 0 to  $\pi$ , there is a continuous transition from one scenario to the other, with an inflection point at  $\frac{\pi}{2}$  where the lightest particle transitions from the neutralino to the chargino.

### 3.5 Wino inelastic magnetic dipole

The wino inelastic dipole with Lagrangian in eq. (2.32) is shown in figure 4. The results and physics are similar to the higgsino case. When the cut-off is large, the dipole component has a negligible effect, and the resultant relic density occurs at a mass equal to the canonical 3 TeV. However, such relic densities are ruled out by indirect detection, magnified due to the Sommerfeld enhancement [70, 82, 83]. Again, where the dipole is large, there is additional annihilation in the early universe which can reduce the overabundance of the dark matter density for larger masses. The additional annihilation allows for thermal relics with masses up to  $\sim 100$  TeV in the case that  $d_3 = 1$ . The addition of the dipole also allows for indirect detection through the annihilation diagrams arising directly from the dipole. There is therefore a region of the parameter space below  $\Lambda \sim 3$  TeV, but above the  $W$  and  $Z$  masses where the model is excluded by the indirect detection constraints.

### 3.6 Wino-Higgs interaction

Figures 5(a) and 5(b) shows the mass spectrum and parameter space of the operators in eqs. (2.30) which shift the mass of the wino. Like in the higgsino cases, for  $\frac{d_1 v^2}{\Lambda} < M_2$ ,



**Figure 5.** Mass contour plots (left) and parameter space (right) for wino-Higgs operators. For a description of the mass contour plots, see the caption of figure 2. For a description of the parameter space plots, see the caption of figure 1.

the mass term is dominated by the mass parameters in the Lagrangian and the effect of the dimension-5 term is negligible. For  $\frac{(d_2+d_3)v^2}{\Lambda} > M_2$ , the contributions from the dimension-5 terms dominate, and the relic density increases for decreasing  $\Lambda$ . Thus, the observed relic density follows the mass contour through the parameter space for  $m_\chi = 3$  TeV. However, unlike in the higgsino case, we have an additive mass correction, rather than a cancelling correction, as well as a contribution to the chargino mass. Hence, the behaviour of the relic density is completely determined by the measurable masses of the dark matter particles.

## 4 Discussion

We have introduced a series of effective field theory operators for composite higgsino-like and wino-like dark matter.

The operators contribute to both annihilation to Higgs bosons and gauge bosons, the latter having the form of an inelastic electromagnetic dipole. We have calculated the relic density for a range of values in the mass-coupling parameter space. We then compared the viability to experimental constraints, including direct and indirect detection experiments and LEP bounds.

Most of the operators considered have a region of the parameter space for sufficiently strong couplings where annihilation via the dimension-5 operators dominates the relic density computation. The additional annihilation reduces the relic density as the depletion of dark matter is more efficient. However, the operators which decay into a pair of Higgs bosons require coupling strengths which are too large to be accounted for by an effective field theory, as the interaction occurs at a scale well above the UV-cutoff of the theory. For the models which involve an electromagnetic dipole, inelastic magnetic dipoles with strength  $\mathcal{O}(10^{-3})$  times the proton magnetic moment result in a viable annihilation pathway within the bounds of the effective field theory. These models allow for viable thermal relic higgsino-like or wino-like dark matter with masses up to  $\sim 10$  TeV or  $\sim 100$  TeV respectively. Higher mass models are not valid in the present effective field theory description. These regions will be detectable through indirect detection experiments with a  $\mathcal{O}(10)$  increase in the sensitivity



of indirect detection experiments. In particular, the proposed sensitivities of the Cherenkov Telescope Array [88] should cover the remaining viable areas of the parameter space for the electromagnetic dipole models, but not necessarily in the regions of viability arising from the cancellation of the mass parameters.

We have also seen that the operators which couple to a Higgs boson can produce shifts in the masses of the neutralinos once the Higgs bosons have undergone electroweak symmetry breaking. The operators produces mass splittings between the neutralinos and between the neutralinos and charginos. These splittings cause the lightest neutralino mass to approach zero due to a cancellation in the mass term between the given mass and the dimension-5 coupling. As the mass decreases, the loss of degeneracy to the charginos reduces the efficiency of the coannihilation processes. There is then a region of the parameter space which is a viable thermal relic, with a masses near 100 GeV. These regions are bound by the LEP constraints from the invisible  $Z$ -width, and may be within range of near-future experiments. The required strength of these operators is too big to be produced by the conventional MSSM, but may be produced by an alternative higher scale theory.

Also, in general, the higher order operators we consider do not appear in isolation. For most of the operators that we have considered, they are the only operator added to the model, except where we have combined operators which perform similar functions. The effect of combining operators which directly annihilate is to make the depletion processes more efficient. There may be a roughly two-fold decrease in the size of the interaction required to produce a viable thermal relic. For the Higgs boson interactions, including multiple interactions of varying strength does not change the resultant physics, as there will still in general be a splitting between the neutralinos and chargino. The exception occurs for the higgsinos (winos) where  $c_4(d_5)$  is negative relative to  $\mu(M_2)$  and in the case of higgsinos is equal to the average of  $c_1$  and  $c_2$ .

Here, we have only considered higgsino-like or wino-like dark matter, where the other SUSY neutralino components (if they exist at all) are assumed to be at a significantly higher scale and are integrated out of the equations. Various models exist where the higgsinos and a wino component occur at the same scale, including the so-called “well-tempered neutralino” [11, 16]. While all of the operators we consider also apply to a general four neutralino model, there are no gauge-invariant combinations which mix a wino and a higgsino, though there are Yukawa terms which mix the neutralinos. The additional mixing will impact the particle masses near the low mass viable region.

The effective field theory approach allows for simplified models of neutralino dark matter to be analysed without specifying the higher order theory. Contrary to results from the MSSM, where the thermal relic higgsino-like and wino-like dark matter is constrained to be one value, viable thermal relic higgsinos exist in a wide spectrum from a few tens of GeV to a few tens of TeV. These regions may be constrained in the near future, but composite thermal relic higgsinos remain a key model in the search for dark matter.

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